

$B^0 \rightarrow K^{*0} \mu^+ \mu^-$ Decay in the Aligned Two-Higgs-Doublet Model

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Abstract

In the aligned two-Higgs-doublet model, we perform a complete one-loop computation of the short-distance Wilson coefficients $C_{7,9,10}^{(\prime)}$, which are the most relevant ones for $b \rightarrow s \ell^+ \ell^-$ transitions. It is found that, when the model parameter $|\varsigma_u|$ is much smaller than $|\varsigma_d|$, the charged-scalar contributes mainly to chirality-flipped $C'_{9,10}$, with the corresponding effects being proportional to $|\varsigma_d|^2$. Numerically, the charged-scalar effects fit into two categories: (A) $C_{7,9,10}^{H^\pm}$ are sizable, but $C'_{9,10} \simeq 0$; (B) $C_7^{H^\pm}$ and $C_{9,10}^{H^\pm}$ are sizable, but $C_{9,10}^{H^\pm} \simeq 0$. Taking into account phenomenological constraints from the inclusive radiative decay $B \rightarrow X_s \gamma$, as well as the latest model-independent global analysis of $b \rightarrow s \ell^+ \ell^-$ data, we obtain the much restricted parameter space of the model. We then study the impact of the allowed model parameters on the angular observables P_2 and P'_5 of $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ decay, and find that P'_5 could be increased significantly to be consistent with the experimental data in case B.

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1 Introduction

The rare $B \rightarrow K^* \ell^+ \ell^-$ decays, being the flavour-changing neutral-current (FCNC) processes, do not arise at tree level and are highly suppressed at higher orders within the Standard Model (SM), due to the Glashow-Iliopoulos-Maiani (GIM) mechanism [1]. However, new TeV-scale particles in many extensions of the SM could affect the decay amplitude at a similar level as the SM does. These decays play, therefore, a crucial role in testing the SM and probing various NP scenarios beyond it [2]. It is particularly interesting to note that, based on these decays, observables with a very limited sensitivity to hadronic uncertainties can be constructed, enhancing therefore the discovery potential for NP [3–6].

Experimentally, several interesting deviations from the SM predictions have been observed in these decays. In 2013, the form-factor-independent angular observable P'_5 [4, 5] of $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ decay was measured by the LHCb collaboration [7], showing a 3.7σ disagreement with the SM expectation [8–11]. Recently, the LHCb collaboration has released new measurements of the angular observables in this decay, based on the dataset of 3 fb^{-1} of integrated luminosity, and still found a 3.4σ deviation for P'_5 [12]. Moreover, being in agreement with the LHCb measurements, a deviation with a significance of 2.1σ was also reported by the Belle collaboration [13]. Besides the P'_5 anomaly, there are some other slight deviations beyond the 2σ level, such as the observables P_2 in $q^2 \in [2, 4.3] \text{ GeV}^2$ and P'_4 in $q^2 \in [14.18, 16] \text{ GeV}^2$ [14–16]. These anomalies have triggered lots of theoretical studies both within the SM and in various NP models [4–6, 8–11, 14–36].

As a minimal extension of the SM scalar sector, the two-Higgs-doublet model (2HDM) [37] can easily satisfy the electroweak (EW) precision data and, at the same time, lead to a very rich phenomenology [38]. The scalar spectrum consists of two charged scalars H^\pm and three neutral ones h , H , and A , one of which is to be identified with the SM-like Higgs boson found at the LHC [39, 40]. The direct search for these additional scalar states would be an important task for high-energy colliders in the next few years. It should be noted that, complementary to the direct searches, indirect constraints on the 2HDM could also be obtained from the rare FCNC decays like $B \rightarrow K^* \ell^+ \ell^-$, since these scalars can affect these processes through the penguin and box diagrams. These studies are also very helpful to gain further insights into the scalar sector of supersymmetry and other models that contain similar scalar contents [41–43].

In a generic 2HDM, the non-diagonal couplings of neutral scalars to fermions lead to tree-

level FCNC interactions, which can be avoided by imposing on the Lagrangian an ad-hoc discrete \mathcal{Z}_2 symmetry. Depending on the \mathcal{Z}_2 charge assignments to the scalars and fermions, this results in four types of 2HDMs (types I, II, X, Y) [38, 44] under the hypothesis of natural flavour conservation (NFC) [45]. In the aligned two-Higgs-doublet model (A2HDM) [46], on the other hand, the absence of tree-level FCNCs is automatically guaranteed by assuming the alignment in flavour space of the Yukawa matrices for each type of right-handed fermions. Interestingly, the A2HDM can recover as particular cases all known specific implementations of the 2HDMs based on \mathcal{Z}_2 symmetries. The model is also featured by possible new sources of CP violation beyond that of the Cabibbo-Kobayashi-Maskawa (CKM) matrix [47, 48]. These features make the A2HDM very attracting both in high-energy collider physics [49–55] and in low-energy flavour physics [56–66].

In this paper, we will study the decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ in the A2HDM. Our paper is organized as follows: In section 2, we give a brief overview of the A2HDM, focusing mainly on the scalar and Yukawa sectors. In section 3, a complete one-loop computation of the short-distance (SD) Wilson coefficients $C_{7,9,10}^{(\prime)}$ is presented, and the final analytical expressions are given both within the SM and in the A2HDM. The angular observables of $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ decay are also introduced in this section. In section 4, taking into account phenomenological constraints from the inclusive radiative decay $B \rightarrow X_s \gamma$ and the latest model-independent global analysis of $b \rightarrow s \ell^+ \ell^-$ data, we study the impact of the allowed model parameters on the angular observables P_2 and P'_5 of $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ decay. Finally, our conclusions are made in section 5. Some relevant functions for the Wilson coefficients are collected in the appendices.

2 The aligned two-Higgs doublet model

We consider the minimal version of 2HDM, which is invariant under the SM gauge group and includes, besides the SM matter and gauge fields, two complex scalar $\text{SU}(2)_L$ doublets,

$$\phi_a^T(x) = \frac{e^{i\theta_a}}{\sqrt{2}} \left(\sqrt{2} \varphi_a^+, v_a + \rho_a + i\eta_a \right), \quad (a = 1, 2), \quad (2.1)$$

with the hypercharge $Y = 1/2$. The neutral components of the two scalar doublets acquire the vacuum expectation values (VEVs) $\langle 0 | \phi_a^T(x) | 0 \rangle = (0, v_a e^{i\theta_a} / \sqrt{2})$. Through an appropriate $\text{U}(1)_Y$ transformation, one can enforce $\theta_1 = 0$ and leave the relative phase $\theta = \theta_2 - \theta_1$ as

physical. Using further a global $SU(2)$ transformation in the scalar space, one can rotate the original scalar basis to the so-called Higgs basis [67–69],

$$\begin{pmatrix} \Phi_1 \\ -\Phi_2 \end{pmatrix} \equiv \begin{pmatrix} \cos \beta & \sin \beta \\ \sin \beta & -\cos \beta \end{pmatrix} \begin{pmatrix} \phi_1 \\ e^{-i\theta_a} \phi_2 \end{pmatrix}, \quad (2.2)$$

where the rotation angle (clockwise) $\tan \beta = v_2/v_1$. In the new basis, only the scalar doublet Φ_1 gets a nonzero VEV $\langle 0 | \Phi_1^T(x) | 0 \rangle = (0, v/\sqrt{2})$, with $v = \sqrt{v_1^2 + v_2^2} = (\sqrt{2}G_F)^{-1/2} \simeq 246$ GeV, and the two scalar doublets are now parametrized, respectively, by [46]

$$\Phi_1 = \begin{pmatrix} G^+ \\ \frac{1}{\sqrt{2}}(v + S_1 + iG^0) \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} H^+ \\ \frac{1}{\sqrt{2}}(S_2 + iS_3) \end{pmatrix}, \quad (2.3)$$

where G^\pm and G^0 denote the massless Goldstone fields to be eaten by the W^\pm and Z^0 gauge bosons, respectively. The remaining five physical degrees of freedom are given by the two charged fields $H^\pm(x)$ and the three neutral ones $\varphi_i^0(x) = \{h(x), H(x), A(x)\} = \mathcal{R}_{ij}S_j$, where \mathcal{R} is an orthogonal matrix obtained after diagonalizing the mass terms in the scalar potential. Generally, none of these three neutral scalars can have a definite CP quantum number.

2.1 Scalar sector

The most general scalar potential for the two doublets Φ_1 and Φ_2 that is allowed by the EW gauge symmetry can be written as [67–69]:

$$\begin{aligned} V = & \mu_1 \left(\Phi_1^\dagger \Phi_1 \right) + \mu_2 \left(\Phi_2^\dagger \Phi_2 \right) + \left[\mu_3 \left(\Phi_1^\dagger \Phi_2 \right) + \mu_3^* \left(\Phi_2^\dagger \Phi_1 \right) \right] \\ & + \lambda_1 \left(\Phi_1^\dagger \Phi_1 \right)^2 + \lambda_2 \left(\Phi_2^\dagger \Phi_2 \right)^2 + \lambda_3 \left(\Phi_1^\dagger \Phi_1 \right) \left(\Phi_2^\dagger \Phi_2 \right) + \lambda_4 \left(\Phi_1^\dagger \Phi_2 \right) \left(\Phi_2^\dagger \Phi_1 \right) \\ & + \left[\left(\lambda_5 \Phi_1^\dagger \Phi_2 + \lambda_6 \Phi_1^\dagger \Phi_1 + \lambda_7 \Phi_2^\dagger \Phi_2 \right) \left(\Phi_1^\dagger \Phi_2 \right) + \text{h.c.} \right]. \end{aligned} \quad (2.4)$$

The hermiticity of the potential requires the parameters $\mu_{1,2}$ and $\lambda_{1,2,3,4}$ to be real, while μ_3 and $\lambda_{5,6,7}$ could be generally complex. The minimization condition imposes the relations $\mu_1 = -\lambda_1 v^2$ and $\mu_3 = -\frac{1}{2} \lambda_6 v^2$. Since only the relative phases among $\lambda_{5,6,7}$ are physical, the scalar potential is finally fully characterized by eleven real parameters, v , μ_2 , $\lambda_{1,2,3,4}$, $|\lambda_{5,6,7}|$, $\arg(\lambda_5 \lambda_6^*)$ and $\arg(\lambda_5 \lambda_7^*)$, four of which can be determined by the scalar masses $M_{H^\pm, h, H, A}$.

Explicitly, inserting Eq. (2.3) into Eq. (2.4) and imposing the minimization condition, one gets $M_{H^\pm}^2 = \mu_2 + \frac{1}{2}\lambda_3 v^2$, and the mass-squared matrix \mathcal{M}^2 of $S_{1,2,3}$ fields in terms of v and λ_i . Using the orthogonal matrix \mathcal{R} , one can then obtain the masses of the three neutral scalars, $\mathcal{R}\mathcal{M}^2\mathcal{R}^T = \text{diag}(M_h^2, M_H^2, M_A^2)$.

In the CP-conserving limit, $\lambda_{5,6,7}$ are all real and the neutral scalars are CP eigenstates. The CP-odd scalar A corresponds to S_3 , with the mass given by $M_A^2 = M_{H^\pm}^2 + v^2(\frac{\lambda_4}{2} - \lambda_5)$, while the two CP-even scalars h and H are orthogonal combinations of S_1 and S_2 ,

$$\begin{pmatrix} h \\ H \end{pmatrix} = \begin{pmatrix} \cos \tilde{\alpha} & \sin \tilde{\alpha} \\ -\sin \tilde{\alpha} & \cos \tilde{\alpha} \end{pmatrix} \begin{pmatrix} S_1 \\ S_2 \end{pmatrix}, \quad (2.5)$$

where the mixing angle $\tilde{\alpha}$ is determined by

$$\tan \tilde{\alpha} = \frac{M_h^2 - 2\lambda_1 v^2}{v^2 \lambda_6} = \frac{v^2 \lambda_6}{2\lambda_1 v^2 - M_H^2}. \quad (2.6)$$

The masses of the two neutral scalars are given, respectively, by $M_h^2 = \frac{1}{2}(\Sigma - \Delta)$ and $M_H^2 = \frac{1}{2}(\Sigma + \Delta)$, where

$$\begin{aligned} \Sigma &= M_{H^\pm}^2 + v^2 \left(2\lambda_1 + \frac{\lambda_4}{2} + \lambda_5 \right), \\ \Delta &= \sqrt{\left[M_{H^\pm}^2 + v^2 \left(-2\lambda_1 + \frac{\lambda_4}{2} + \lambda_5 \right) \right]^2 + 4v^4(\lambda_6)^2} = -\frac{2v^2 \lambda_6}{\sin(2\tilde{\alpha})}. \end{aligned} \quad (2.7)$$

Here $M_h \leq M_H$ by convention and the SM limit is recovered when $\tilde{\alpha} = 0$.

2.2 Yukawa sector

The Yukawa Lagrangian of the 2HDM is most generally given by [38, 46]

$$\mathcal{L}_Y = - \left[\bar{Q}'_L (\Gamma_1 \phi_1 + \Gamma_2 \phi_2) d'_R + \bar{Q}'_L (\Delta_1 \tilde{\phi}_1 + \Delta_2 \tilde{\phi}_2) u'_R + \bar{L}'_L (\Pi_1 \phi_1 + \Pi_2 \phi_2) \ell'_R \right] + \text{h.c.}, \quad (2.8)$$

where $\tilde{\phi}_a(x) \equiv i\tau_2 \phi_a^*(x)$ are the charge-conjugated fields with $Y = -\frac{1}{2}$, \bar{Q}'_L and \bar{L}'_L are the left-handed quark and lepton doublets, and u'_R , d'_R and ℓ'_R the corresponding right-handed singlets, in the weak-interaction basis. All fermionic fields are written as 3-dimensional vectors and the couplings Γ_a , Δ_a and Π_a are therefore 3×3 complex matrices in flavour space.

Transforming to the Higgs basis, Eq. (2.8) becomes

$$\mathcal{L}_Y = -\frac{\sqrt{2}}{v} \left[\bar{Q}'_L (M'_d \Phi_1 + Y'_d \Phi_2) d'_R + \bar{Q}'_L (M'_u \tilde{\Phi}_1 + Y'_u \tilde{\Phi}_2) u'_R + \bar{L}'_L (M'_\ell \Phi_1 + Y'_\ell \Phi_2) \ell'_R \right] + \text{h.c.}, \quad (2.9)$$

where

$$M'_d = \frac{1}{\sqrt{2}} (v_1 \Gamma_1 + v_2 \Gamma_2 e^{i\theta}), \quad Y'_d = \frac{1}{\sqrt{2}} (-v_2 \Gamma_1 + v_1 \Gamma_2 e^{i\theta}), \quad (2.10)$$

$$M'_u = \frac{1}{\sqrt{2}} (v_1 \Delta_1 + v_2 \Delta_2 e^{-i\theta}), \quad Y'_u = \frac{1}{\sqrt{2}} (-v_2 \Delta_1 + v_1 \Delta_2 e^{-i\theta}), \quad (2.11)$$

$$M'_\ell = \frac{1}{\sqrt{2}} (v_1 \Pi_1 + v_2 \Pi_2 e^{i\theta}), \quad Y'_\ell = \frac{1}{\sqrt{2}} (-v_2 \Pi_1 + v_1 \Pi_2 e^{i\theta}). \quad (2.12)$$

In general, the Yukawa matrices M'_f and Y'_f ($f = u, d, \ell$) cannot be simultaneously diagonalized in flavour space. Thus, in the mass-eigenstate basis, with diagonal fermion mass matrices M_f , the corresponding Yukawa matrices Y_f remain non-diagonal, giving rise to tree-level FCNC interactions. The unwanted tree-level FCNCs can be eliminated by requiring the alignment in flavour space of the Yukawa matrices [46]:

$$\begin{aligned} \Gamma_2 &= \xi_d e^{-i\theta} \Gamma_1, & \Delta_2 &= \xi_u^* e^{i\theta} \Delta_1, & \Pi_2 &= \xi_\ell e^{-i\theta} \Pi_1, \\ Y_{d,\ell} &= \varsigma_{d,\ell} M_{d,\ell}, & Y_u &= \varsigma_u^* M_u, & \varsigma_f &\equiv \frac{\xi_f - \tan \beta}{1 + \xi_f \tan \beta}, \end{aligned} \quad (2.13)$$

where ξ_f (ς_f) are arbitrary complex parameters and could introduce new sources of CP violation beyond that of the CKM matrix. The interactions of the charged scalar with the fermion mass-eigenstate fields then read

$$\mathcal{L}_{H^\pm} = -\frac{\sqrt{2}}{v} H^+ \left\{ \bar{u} \left[\varsigma_d V_{\text{CKM}} M_d P_R - \varsigma_u M_u^\dagger V_{\text{CKM}} P_L \right] d + \varsigma_\ell \bar{\nu} M_\ell P_R \ell \right\} + \text{h.c.}, \quad (2.14)$$

where $P_{L(R)} \equiv (1 \mp \gamma_5)/2$ is the left (right)-handed chirality projector, and V_{CKM} the CKM matrix [47, 48]. Here we did not give the neutral scalar sector [46] in \mathcal{L}_Y or the FCNC local structures induced beyond tree-level (quantum corrections) [56], because their effects are highly suppressed by the muon mass in the decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$. The usual NFC models [38, 44], with discrete \mathcal{Z}_2 symmetries, are recovered for particular values of ς_f , as shown in Table 1.

Table 1: The one-to-one correspondence between different specific choices of the couplings ς_f and the 2HDMs based on discrete \mathcal{Z}_2 symmetries.

Model	ς_d	ς_u	ς_l
Type I	$\cot \beta$	$\cot \beta$	$\cot \beta$
Type II	$-\tan \beta$	$\cot \beta$	$-\tan \beta$
Type X	$\cot \beta$	$\cot \beta$	$-\tan \beta$
Type Y	$-\tan \beta$	$\cot \beta$	$\cot \beta$
Inert	0	0	0

3 $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ in the A2HDM

3.1 Effective weak Hamiltonian

The rare decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ proceeds through the loop diagrams both within the SM and in the A2HDM. When the heavy degrees of freedom, including the top quark, the weak gauge bosons, as well as the charged scalars, have been integrated out, we obtain the low-energy effective weak Hamiltonian governing the decay [3, 70]:

$$\mathcal{H}_{\text{eff}} = -\frac{4G_F}{\sqrt{2}} V_{tb} V_{ts}^* \sum_i (C_i O_i + C'_i O'_i) , \quad (3.1)$$

where G_F is the Fermi coupling constant. Here we neglect the doubly Cabibbo-suppressed (proportional to $V_{ub} V_{us}^*$) contributions to Eq. (3.1), and focus only on the operators [3]:

$$O_7 = \frac{e}{16\pi^2} \bar{m}_b (\bar{s} \sigma^{\mu\nu} P_R b) F_{\mu\nu} , \quad O'_7 = \frac{e}{16\pi^2} \bar{m}_b (\bar{s} \sigma^{\mu\nu} P_L b) F_{\mu\nu} , \quad (3.2)$$

$$O_9 = \frac{e^2}{16\pi^2} (\bar{s} \gamma^\mu P_L b) (\bar{\mu} \gamma_\mu \mu) , \quad O'_9 = \frac{e^2}{16\pi^2} (\bar{s} \gamma^\mu P_R b) (\bar{\mu} \gamma_\mu \mu) , \quad (3.3)$$

$$O_{10} = \frac{e^2}{16\pi^2} (\bar{s} \gamma^\mu R_L b) (\bar{\mu} \gamma_\mu \gamma_5 \mu) , \quad O'_{10} = \frac{e^2}{16\pi^2} (\bar{s} \gamma^\mu P_R b) (\bar{\mu} \gamma_\mu \gamma_5 \mu) , \quad (3.4)$$

where $\bar{m}_b = \bar{m}_b(\mu)$ denotes the b -quark running mass in the $\overline{\text{MS}}$ scheme.

Within the SM, the electromagnetic dipole operator O_7 and the semileptonic operators $O_{9,10}$ play the leading role in the decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$. Besides modifying the values of the SD Wilson coefficients $C_{7,9,10}$, the charged-scalar contributions could also make the chirality-flipped operators $O'_{7,9,10}$ defined above to contribute in a significant manner, especially in some regions

of the parameter spaces discussed later.

The SD Wilson coefficients $C_i(\mu)$ and $C'_i(\mu)$ can be obtained firstly at the matching scale $\mu_W \sim M_W$ perturbatively, by requiring equality of the one-particle irreducible Green functions calculated in the full and in the effective theory [70]. Using the renormalization group equation, one can then get $C_i(\mu)$ and $C'_i(\mu)$ at the lower scale $\mu_b \sim m_b$. During the calculation, the limit $\bar{m}_{u,c} \rightarrow 0$ and the unitarity of the CKM matrix have been used. For simplicity, we introduce the mass ratios:

$$x_t = \frac{\bar{m}_t^2(\mu_W)}{M_W^2}, \quad y_t = \frac{\bar{m}_t^2(\mu_W)}{M_{H^\pm}^2}. \quad (3.5)$$

Details of the computational method could be found, for example, in refs. [62, 70].

3.2 Wilson coefficients in the SM

In the SM, the one-loop penguin and box diagrams have been calculated both in the Feynman ($\xi = 1$) and in the unitary ($\xi = \infty$) gauge [71–79], denoted by the subscript ‘F’ and ‘U’, respectively. The different contributions to $C_i^{\text{SM}}(\mu_W)$ can be split into the following forms:

$$C_7^{\text{SM}} = C_7^{\gamma\text{-penguin}}, \quad (3.6)$$

$$C_9^{\text{SM}} = C_9^{W\text{-box}} + C_9^{Z\text{-penguin}} + C_9^{\gamma\text{-penguin}}, \quad (3.7)$$

$$C_{10}^{\text{SM}} = C_{10}^{W\text{-box}} + C_{10}^{Z\text{-penguin}}, \quad (3.8)$$

where the corresponding parts resulting from the W -box, Z -penguin and γ -penguin diagrams are given, respectively, by

$$C_{9,\text{F(U)}}^{W\text{-box}} = -\frac{B_{0,\text{F(U)}}}{\sin^2 \theta_W}, \quad C_{10,\text{F(U)}}^{W\text{-box}} = \frac{B_{0,\text{F(U)}}}{\sin^2 \theta_W}, \quad (3.9)$$

$$C_{9,\text{F(U)}}^{Z\text{-penguin}} = \left(-4 + \frac{1}{\sin^2 \theta_W}\right) C_{0,\text{F(U)}}, \quad C_{10,\text{F(U)}}^{Z\text{-penguin}} = -\frac{C_{0,\text{F(U)}}}{\sin^2 \theta_W}, \quad (3.10)$$

$$C_{7,\text{F(U)}}^{\gamma\text{-penguin}} = -\frac{1}{2} D'_{0,\text{F(U)}}, \quad C_{9,\text{F(U)}}^{\gamma\text{-penguin}} = -D_{0,\text{F(U)}} + \frac{4}{9}, \quad (3.11)$$

where θ_W is the weak mixing angle, and the Inami-Lim functions [71] are defined as

$$B_{0,\text{F}} = F_1(x_t), \quad C_{0,\text{F}} = F_3(x_t), \quad D'_{0,\text{F}} = F_6(x_t), \quad D_{0,\text{F}} = -\frac{4}{9}F_0(x_t) + F_5(x_t), \quad (3.12)$$

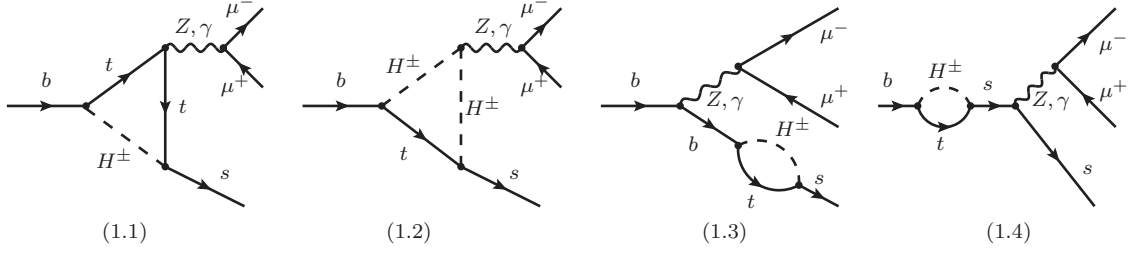


Figure 1: Z - and γ -penguin diagrams involving the charged-scalar exchanges in the A2HDM.

in the Feynman gauge, and

$$\begin{aligned}
B_{0,U} &= -\frac{x_t}{16}L_\epsilon + F_4(x_t), & C_{0,U} &= -\frac{x_t}{16}L_\epsilon - F_1(x_t) + F_3(x_t) + F_4(x_t), \\
D'_{0,U} &= F_6(x_t), & D_{0,U} &= \frac{x_t}{4}L_\epsilon - \frac{4}{9}F_0(x_t) + 4F_1(x_t) - 4F_4(x_t) + F_5(x_t), \quad (3.13)
\end{aligned}$$

in the unitary gauge. Here we introduce the notation $L_\epsilon \equiv \frac{1}{\epsilon} + \log\left(\frac{\mu_W^2}{M_W^2}\right)$, where $\epsilon = (4-d)/2$ is the dimensional regulator of ultraviolet divergence. Explicit expressions of the basic functions $F_i(x)$ are given by Eqs. (A.1)–(A.9). While each piece on the right-hand side of Eqs. (3.7) and (3.8) depends obviously on ϵ in the unitary gauge, due to the longitudinal components of the W^\pm , Z^0 and off-shell photon propagators, the physical quantities $C_{7,9,10}^{\text{SM}}$ are indeed free of ϵ and are independent of the EW gauge fixings. For a recent review of higher-order corrections to $C_{7,9,10}^{\text{SM}}$, the readers are referred to ref. [80].

3.3 Wilson coefficients in the A2HDM

In the A2HDM, the charged-scalar exchanges lead to additional contributions to $C_{7,9,10}$ and could also make the chirality-flipped operators $O'_{7,9,10}$ to contribute in a significant manner, through the Z^0 - and γ -penguin diagrams shown in Figure 1. Since we have neglected the light lepton mass, there is no contribution from the SM W -box diagrams with the W^\pm bosons replaced by the charged scalars H^\pm .

For each Feynman diagram shown in Figure 1, the contributions are identical in the two gauges. The total Wilson coefficients $C_{7,9,10}$ are split into two parts, one is from the SM contributions $C_{7,9,10}^{\text{SM}}$, and the other from the charged-scalar ones $C_{7,9,10}^{\text{H}^\pm}$. For the chirality-flipped operators, $C'_{7,9,10} = C_{7,9,10}^{\text{H}^\pm}$, because the SM contributions are well suppressed by the factor \bar{m}_s/\bar{m}_b . For convenience, we decompose these new contributions in such a way to render

explicit their dependence on the couplings ς_u and ς_d :

$$C_7^{\text{H}\pm} = |\varsigma_u|^2 C_{7,\text{uu}} + \varsigma_d \varsigma_u^* C_{7,\text{ud}}, \quad (3.14)$$

$$C_9^{\text{H}\pm} = |\varsigma_u|^2 C_{9,\text{uu}}, \quad (3.15)$$

$$C_{10}^{\text{H}\pm} = |\varsigma_u|^2 C_{10,\text{uu}}, \quad (3.16)$$

$$C_7^{\prime\text{H}\pm} = \frac{\bar{m}_s}{\bar{m}_b} (|\varsigma_u|^2 C_{7,\text{uu}} + \varsigma_u \varsigma_d^* C_{7,\text{ud}}), \quad (3.17)$$

$$C_9^{\prime\text{H}\pm} = (-1 + 4 \sin^2 \theta_W) C_{10}^{\prime\text{H}\pm} + \frac{\bar{m}_b \bar{m}_s}{M_W^2} \left[|\varsigma_u|^2 C'_{9,\text{uu}} + 2\Re(\varsigma_u \varsigma_d^*) C'_{9,\text{ud}} + |\varsigma_d|^2 C'_{9,\text{dd}} \right], \quad (3.18)$$

$$C_{10}^{\prime\text{H}\pm} = \frac{\bar{m}_b \bar{m}_s}{M_W^2} \left[|\varsigma_u|^2 C'_{10,\text{uu}} + 2\Re(\varsigma_u \varsigma_d^*) C'_{10,\text{ud}} + |\varsigma_d|^2 C'_{10,\text{dd}} \right], \quad (3.19)$$

where the coefficients of the different combinations of the couplings ς_u and ς_d are given by Eqs. (B.1)–(B.10). In the particular cases of type II and type Y 2HDMs with large $\tan \beta$, the only terms enhanced by a factor $\tan^2 \beta$ originate from the $|\varsigma_d|^2$ part contributing only to $C_{9,10}^{\prime\text{H}\pm}$. The Wilson coefficients $C_{7,9,10}^{(\prime)\text{H}\pm}$ are found to be invariant under the transformations $\varsigma_u \rightarrow -\varsigma_u$ and $\varsigma_d \rightarrow -\varsigma_d$. There is an implicit μ_W dependence via the s, b, t -quark masses, which depend on the precise definitions and have to be specified when going beyond the leading logarithm (LL). As we evaluate $C_{7,9,10}^{(\prime)\text{H}\pm}$ only at the leading order (LO) in α_s , whether the running masses $\bar{m}_q(\mu_W)$ or the pole masses m_q are used does not matter too much. As a consequence, we choose the pole masses m_q as input in Eqs. (3.17)–(3.19).

Our results for the chirality-flipped Wilson coefficients $C_{7,9,10}^{\prime\text{H}\pm}$ are presented for the first time in the A2HDM. In the particular cases of the \mathcal{Z}_2 symmetric 2HDMs, our results agree with the ones calculated in refs. [81–84]. It is also noted that the next-to-leading order QCD corrections to $C_{7,9,10}^{\text{H}\pm}$ in the supersymmetry and type-II 2HDM have already been calculated in refs. [85–89].

3.4 Angular observables in $B^0 \rightarrow K^{*0} \mu^+ \mu^-$ decay

The angular distribution of the $B^0 \rightarrow K^{*0}(\rightarrow K^+ \pi^-) \mu^+ \mu^-$ decay is described by the dimuon invariant mass squared q^2 as well as the three angles θ_ℓ , θ_{K^*} and ϕ , where θ_ℓ is defined as the angle between the flight direction of the μ^+ (μ^-) and the opposite direction of the B^0 (\bar{B}^0) in the rest frame of the dimuon system, and θ_{K^*} the angle between the flight direction of the K^+ (K^-) and that of the B^0 (\bar{B}^0) in the K^{*0} (\bar{K}^{*0}) rest frame, while ϕ is the angle between the

plane containing the dimuon pair and the plane containing K^+ and π^- mesons in the B^0 (\bar{B}^0) rest frame. In terms of these four kinematic variables, the full angular decay distribution of the decay is then given by [3, 90]

$$\begin{aligned} \frac{d^4\bar{\Gamma}[B^0 \rightarrow K^{*0}\mu^+\mu^-]}{dq^2 d\cos\theta_\ell d\cos\theta_{K^*} d\phi} = \frac{9}{32\pi} & \left[\bar{I}_1^s \sin^2\theta_{K^*} + \bar{I}_1^c \cos^2\theta_{K^*} + (\bar{I}_2^s \sin^2\theta_{K^*} + \bar{I}_2^c \cos^2\theta_{K^*}) \cos 2\theta_\ell \right. \\ & + \bar{I}_3 \sin^2\theta_{K^*} \sin^2\theta_\ell \cos 2\phi + \bar{I}_4 \sin 2\theta_{K^*} \sin 2\theta_\ell \cos \phi \\ & + \bar{I}_5 \sin 2\theta_{K^*} \sin \theta_\ell \cos \phi \\ & + \bar{I}_6^s \sin^2\theta_{K^*} \cos \theta_\ell + \bar{I}_7 \sin 2\theta_{K^*} \sin \theta_\ell \sin \phi \\ & \left. + \bar{I}_8 \sin 2\theta_{K^*} \sin 2\theta_\ell \sin \phi + \bar{I}_9 \sin^2\theta_{K^*} \sin^2\theta_\ell \sin 2\phi \right], \end{aligned} \quad (3.20)$$

where the angular coefficients $\bar{I}_i^{(a)}$ are functions of q^2 only, and the relations $\bar{I}_1^s = 3\bar{I}_2^s$, $\bar{I}_1^c = -\bar{I}_2^c$ and $\bar{I}_6^c = 0$ hold when the muon mass is neglected. The corresponding expression for the CP-conjugated mode $\bar{B}^0 \rightarrow \bar{K}^{*0}(\rightarrow K^-\pi^+)\mu^+\mu^-$ is obtained from Eq. (3.20) by the replacements $\bar{I}_i^{(a)} \rightarrow I_i^{(a)}$ [3, 90]. Explicit forms of the angular coefficients $\bar{I}_i^{(a)}$ ($I_i^{(a)}$) could be found, for example, in refs. [3, 6, 12].

The self-tagging property of the decay $B^0 \rightarrow K^{*0}\mu^+\mu^-$ makes it possible to determine both the CP-averaged and the CP-asymmetric quantities defined, respectively, by [3]

$$S_i^{(a)} = \left(I_i^{(a)} + \bar{I}_i^{(a)} \right) / \left(\frac{d\Gamma}{dq^2} + \frac{d\bar{\Gamma}}{dq^2} \right), \quad A_i^{(a)} = \left(I_i^{(a)} - \bar{I}_i^{(a)} \right) / \left(\frac{d\Gamma}{dq^2} + \frac{d\bar{\Gamma}}{dq^2} \right). \quad (3.21)$$

The previously studied observables, such as the q^2 distributions of the forward-backward asymmetry A_{FB} and the CP asymmetry A_{CP} , can be expressed in terms of these angular observables.

With the structure of the amplitudes at large recoil, it is possible to build clean observables whose sensitivity to the $B \rightarrow K^*$ transition form factors is suppressed by α_s or Λ_{QCD}/m_b [5]. These include the so-called P_i and P'_i observables defined, respectively, by [5, 91, 92]

$$P_1 = \frac{S_3}{2S_2^s}, \quad P_2 = \frac{S_6^s}{8S_2^s}, \quad P_3 = \frac{S_9}{4S_2^s}, \quad (3.22)$$

and

$$P'_4 = \frac{S_4}{2\sqrt{-S_2^s S_2^c}}, \quad P'_5 = \frac{S_5}{2\sqrt{-S_2^s S_2^c}}, \quad P'_6 = \frac{S_7}{2\sqrt{-S_2^s S_2^c}}, \quad P'_8 = \frac{S_8}{2\sqrt{-S_2^s S_2^c}}. \quad (3.23)$$

The numerical impact of charged-scalar contributions to some of these observables will be discussed in the next section.

4 Numerical results and discussions

4.1 Choice of the model parameters

For the considered decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$, only three model parameters, the charged-scalar mass M_{H^\pm} and the two alignment parameters ς_u and ς_d , are involved. In the following we assume the parameters $\varsigma_{u,d}$ to be real, indicating that the only source of CP violation in the A2HDM is still due to the CKM matrix. Following the previous studies, we give below the preset ranges of these model parameters:

- The charged-scalar mass is assumed to lie in the range $M_{H^\pm} \in [80, 1000]$ GeV, where the lower bound comes from the LEP direct search [93], while the upper bound from the unitarity and stability of the scalar potential [94–96].
- The alignment parameter ς_u is assumed to lie in the range $|\varsigma_u| \leq 2$, to be compatible with the current data of loop-induced processes, such as $Z \rightarrow b\bar{b}$, $b \rightarrow s\gamma$, $B_{s,d}^0 - \bar{B}_{s,d}^0$ mixings, as well as the $h(125)$ decays [54, 55, 57–61].
- The alignment parameter ς_d is only mildly constrained through phenomenological requirements that involve additionally other model parameters. So we let it to be a free parameter.
- In the 2HDMs with discrete \mathcal{Z}_2 symmetries, the parameters ς_u and ς_d are not independent but are related to each other through the ratio of the VEVs $\tan\beta = v_2/v_1$. The upper limit for $\tan\beta$ also comes from the unitarity and stability of the scalar potential [94–96]; we assume here $\tan\beta \leq 50$.

4.2 Constraints on the model parameters

For the other input parameters, we take $M_Z = 91.1876$ GeV, $M_W = 80.385$ GeV, $m_t = (174.2 \pm 1.4)$ GeV, $m_b = (4.78 \pm 0.06)$ GeV, and $\bar{m}_s(2 \text{ GeV}) = (96_{-4}^{+8})$ MeV [97]. Since $C_7^{\text{H}^\pm} = \bar{m}_s/\bar{m}_b C_7^{\text{H}^\pm}$ and $\bar{m}_s \ll \bar{m}_b$, the contribution from O_7' will be safely neglected.

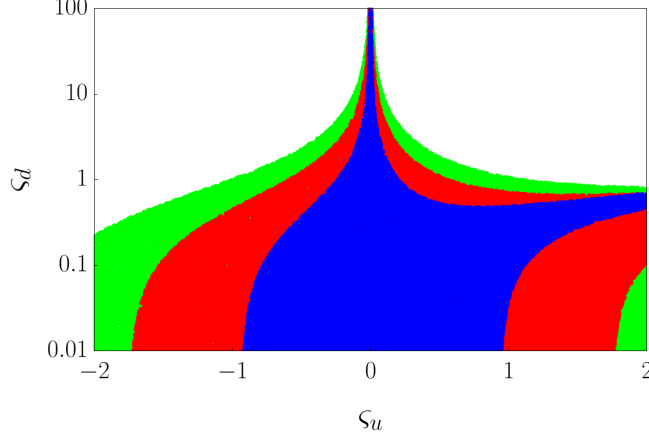


Figure 2: The allowed regions in the $\varsigma_u - \varsigma_d$ plane ($\varsigma_d > 0$) under the constraint from Eq. (4.3). The blue-, red-, and green-band correspond to $M_{H^\pm} = 80, 300$ and 500 GeV, respectively.

The Wilson coefficient $C_7^{\text{H}^\pm}$ is severely constrained by the inclusive decay $B \rightarrow X_s \gamma$. The branching ratio of $B \rightarrow X_s \gamma$ measured by CLEO [98], Belle [99–101] and BaBar [102–104], lead to the combined average [105]

$$\mathcal{B}^{\text{exp}}(B \rightarrow X_s \gamma) \big|_{E_\gamma > 1.6 \text{ GeV}} = (3.43 \pm 0.21_{\text{stat.}} \pm 0.07_{\text{syst.}}) \times 10^{-4}, \quad (4.1)$$

which is in good agreement with the updated SM prediction [106]

$$\mathcal{B}^{\text{SM}}(B \rightarrow X_s \gamma) \big|_{E_\gamma > 1.6 \text{ GeV}} = (3.36 \pm 0.23) \times 10^{-4}. \quad (4.2)$$

It should be noted that the chromomagnetic dipole operator $O_8 = \frac{g_s}{16\pi^2} \bar{m}_b (\bar{s} \sigma^{\mu\nu} P_R T^a b) G_{\mu\nu}^a$ also plays an important role in the decay $B \rightarrow X_s \gamma$. However, at the LO in α_s , this operator contributes to $B \rightarrow X_s \gamma$ only via its mixing with O_7 . It is then found that, at the matching scale $\mu_W = 160$ GeV, the Wilson coefficients $C_7^{\text{H}^\pm}$ and $C_8^{\text{H}^\pm}$ should fulfill the constraint [106]:

$$-0.0634 \leq C_7^{\text{H}^\pm}(\mu_W) + 0.242 C_8^{\text{H}^\pm}(\mu_W) \leq 0.0464, \quad (4.3)$$

where $C_8^{\text{H}^\pm} = |\varsigma_u|^2 C_{8,\text{uu}} + \varsigma_d \varsigma_u^* C_{8,\text{ud}}$ [81], with the functions $C_{8,\text{uu}}$ and $C_{8,\text{ud}}$ given, respectively, by Eqs. (B.11) and (B.12).

Under the constraint from Eq. (4.3), we show in Figure 2 the allowed regions in the $\varsigma_u - \varsigma_d$ plane ($\varsigma_d > 0$), with three representative values of the charged-scalar masses, $M_{H^\pm} = 80, 300$ and 500 GeV as benchmarks. The case with $\varsigma_d < 0$ is obtained from Figure 2 with the

transformations $\varsigma_u \rightarrow -\varsigma_u$ and $\varsigma_d \rightarrow -\varsigma_d$. The allowed range of ς_u (ς_d) becomes quite large when ς_d (ς_u) equals approximately to zero. Particularly, when $\varsigma_u = 0$, it is found that ς_d can vary within a large range. For heavier charged scalars the constraint becomes weaker, because the NP effect starts to decouple.

Motivated by the latest LHCb and Belle measurements of $b \rightarrow s\ell^+\ell^-$ decays, there exist several global fits for the NP contributions to the Wilson coefficients $C_{9,10}^{(\prime)}$ [14–16, 25]. We use two of these global fit results to further constrain the A2HDM parameters. One is obtained from the combined fit to the $b \rightarrow s\ell^+\ell^-$ mesonic decays (at $\mu_b = 4.8$ GeV) [15]:

$$\begin{aligned} -2.2 \leq C_9^{\text{NP}} \leq -0.4, & \quad -0.5 \leq C_{10}^{\text{NP}} \leq 2.0, \\ -1.3 \leq C_9^{\prime\text{NP}} \leq 3.7, & \quad -1.0 \leq C_{10}^{\prime\text{NP}} \leq 1.6, \end{aligned} \quad (4.4)$$

given at the 3σ level. The other includes also the available data on $\Lambda_b \rightarrow \Lambda\mu^+\mu^-$ decay and get (at $\mu_b = 4.2$ GeV) [25]:

$$\begin{aligned} 0.9 \leq C_9^{\text{NP}} \leq 2.5, & \quad 1.8 \leq C_{10}^{\text{NP}} \leq 4.2, \\ -1.3 \leq C_9^{\prime\text{NP}} \leq 1.8, & \quad 1.0 \leq C_{10}^{\prime\text{NP}} \leq 3.1, \end{aligned} \quad (4.5)$$

given at the 1σ level. It is interesting to note that the latter prefers a shift to C_9 that is opposite in sign compared to the former [25]. Since the Wilson coefficients $C_{9,10}^{\text{H}\pm}(\mu_W)$ and $C_{9,10}^{\prime\text{H}\pm}(\mu_W)$ are calculated only at the LO, they should be evolved to the lower scale μ_b at the LL approximation, which means that they are actually not running [107]. Thus, we can apply directly the bounds given by Eqs. (4.4) and (4.5) to $C_{9,10}^{\text{H}\pm}$ and $C_{9,10}^{\prime\text{H}\pm}$. To be more conservative, we require each of these coefficients to lie within the smaller lower and bigger upper bounds of these two global fits. Using these bounds as well as the constraint from Eq. (4.3), we find that the allowed parameter space in the $\varsigma_u - \varsigma_d$ plane are significantly reduced, especially for the model parameter ς_u , as shown in Figure 3. This means that $C_{9,10}^{\text{H}\pm}$ play a major role in the small $|\varsigma_d|$ region ($|\varsigma_d| < 1$) and $C_{9,10}^{\prime\text{H}\pm}$ are quite sizable when ς_u approaches to zero.

It is also interesting to note that, under the constraint from Eq. (4.3) as well as the bounds on $C_{9,10}^{\text{H}\pm}$ and $C_{9,10}^{\prime\text{H}\pm}$ from Eqs. (4.4) and (4.5), we could obtain a bound on ς_d even when ς_u equals to zero. For illustration, the allowed regions in the $\varsigma_d - M_{H^\pm}$ plane when $\varsigma_u = 0$ and in the $\varsigma_u - M_{H^\pm}$

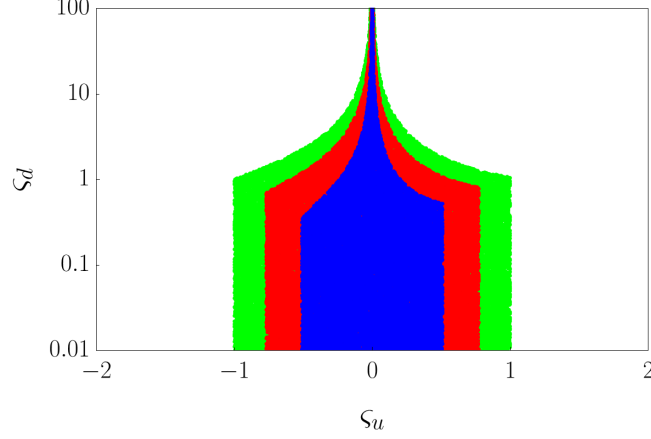


Figure 3: The allowed regions in the $\varsigma_u - \varsigma_d$ plane ($\varsigma_d > 0$) under the constraint from Eq. (4.3) as well as the bounds on $C_{9,10}^{\text{H}\pm}$ and $C_{9,10}'^{\text{H}\pm}$ from Eqs. (4.4) and (4.5). The other captions are the same as in Figure 2.

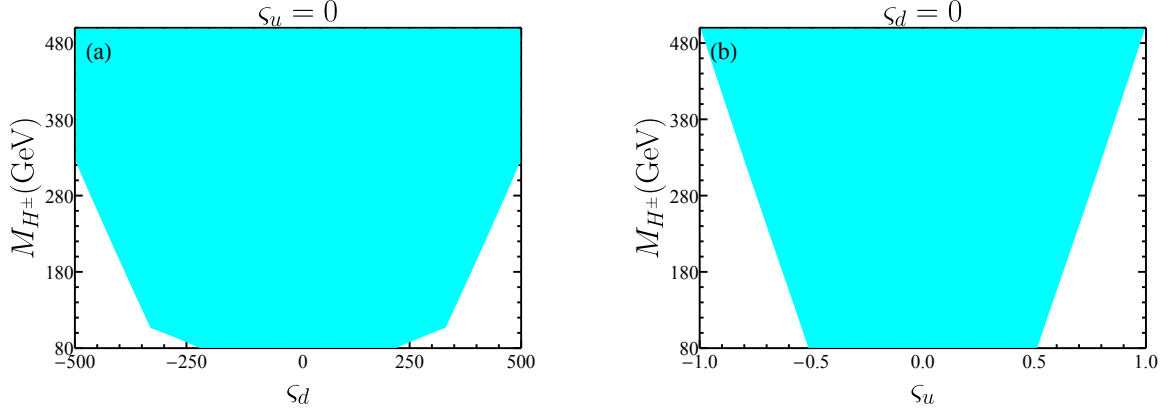


Figure 4: The allowed regions in the $\varsigma_d - M_{H^\pm}$ plane when $\varsigma_u = 0$ (a) and in the $\varsigma_u - M_{H^\pm}$ plane when $\varsigma_d = 0$ (b), under the constraint from Eq. (4.3) as well as the bounds on $C_{9,10}^{\text{H}\pm}$ and $C_{9,10}'^{\text{H}\pm}$ from Eqs. (4.4) and (4.5).

plane when $\varsigma_d = 0$ are shown in Figure 4. Numerically, we obtain $|\varsigma_u| \leq 0.506, 0.763$ and 0.990 , and $|\varsigma_d| \leq 212, 476$ and 622 , corresponding to $M_{H^\pm} = 80, 300$ and 500 GeV, respectively. This means that the more accurate $C_{9,10}'^{\text{NP}}$ can be better used to restrict the parameter ς_d .

4.3 P_2 and P_5' in the A2HDM

In this subsection, with the constrained parameter space, we investigate the impact of A2HDM on the angular observables P_2 and P_5' in the decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$. As there involve only three model parameters in Eqs. (3.14)–(3.19), the five Wilson coefficients ($C_7^{\text{H}\pm}$ is neglected because $\bar{m}_s \ll \bar{m}_b$) are highly correlated with each other, as shown in Figure 5. One can see that, while $C_7^{\text{H}\pm}$ is hardly correlated with the other four Wilson coefficients (Figures 5(a)–5(d)), $C_9^{\text{H}\pm}$

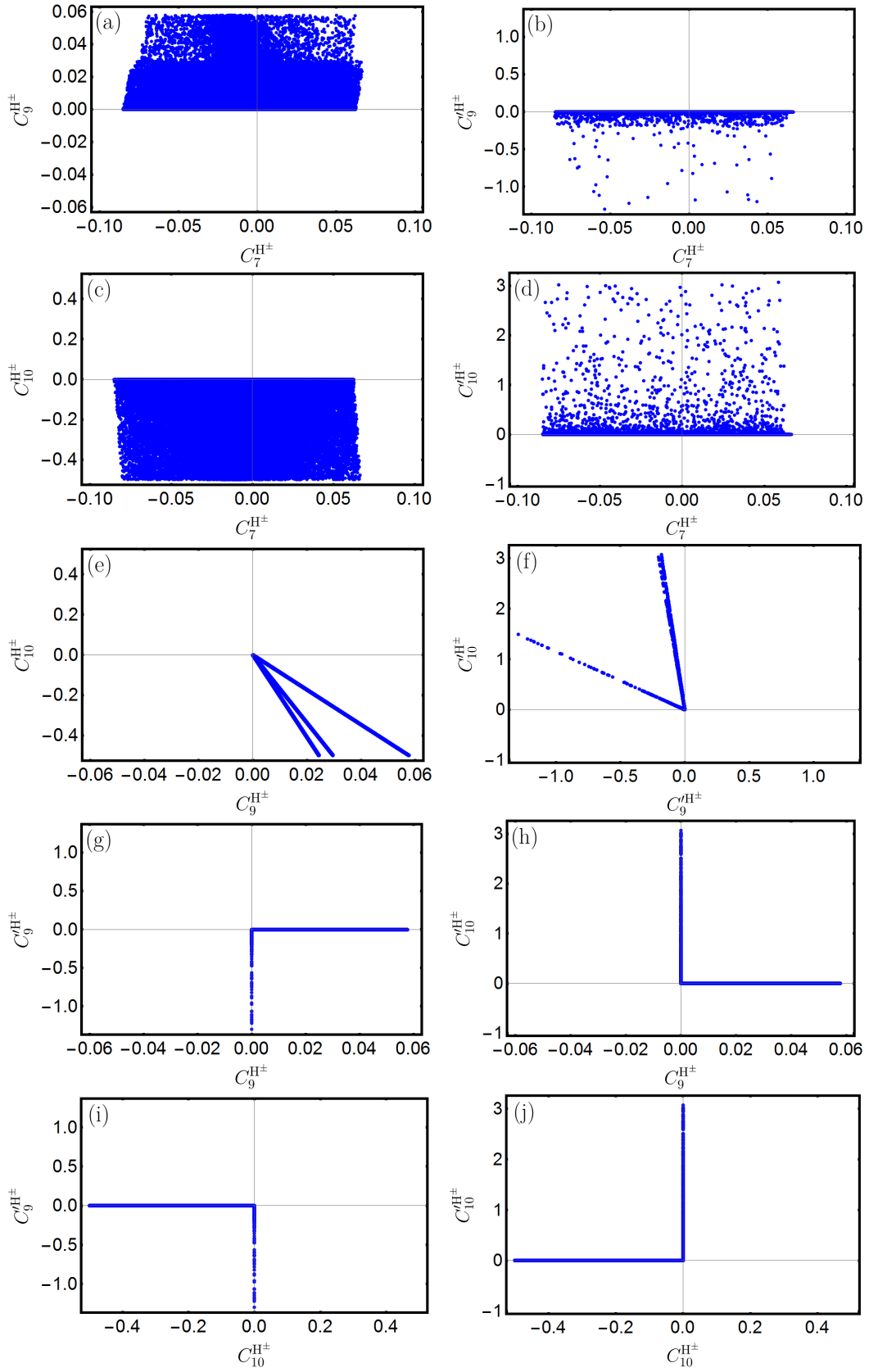


Figure 5: Correlations among the five Wilson coefficients with the allowed ς_u and ς_d values discussed in the previous subsection.

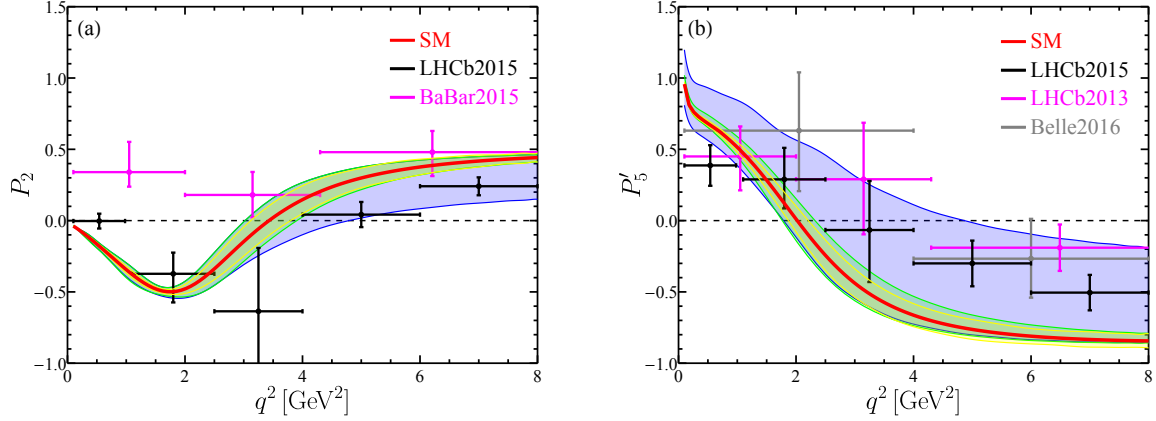


Figure 6: The q^2 dependence of the angular observables P_2 and P'_5 , both within the SM (central value by red and its uncertainty by yellow band) and in the A2HDM (the green and blue bands correspond to the case A and case B, respectively). The experimental data from the LHCb [7, 12], Belle [13] and BaBar [108] collaborations are represented by the corresponding error bars in different q^2 bins.

and $C_{10}^{\text{H}\pm}$ are obviously linearly correlated and the slope depends only on the charged-scalar mass M_{H^\pm} (Figure 5(e)). In addition, $C_9^{\text{H}\pm}$ and $C_{10}^{\text{H}\pm}$ are approximately linearly correlated (Figure 5(f)). The most interesting results are shown in Figures 5(g)–5(j), which suggest that the charged-scalar can not affect the left- and right-handed semileptonic operators at the same time, due to the constraints shown in Figures 2 and 3. These observations motivate us to consider the following two specific cases for the NP Wilson coefficients:

$$\text{Case A: } C_{7,9,10}^{\text{H}\pm} \text{ are sizable, but } C_{9,10}'^{\text{H}\pm} \simeq 0; \quad (4.6)$$

$$\text{Case B: } C_7^{\text{H}\pm} \text{ and } C_{9,10}'^{\text{H}\pm} \text{ are sizable, but } C_{9,10}^{\text{H}\pm} \simeq 0. \quad (4.7)$$

In Figure 6, we show our predictions for the two angular observables P_2 and P'_5 at large K^{*0} recoil both within the SM and in the A2HDM, with the Wilson coefficients obtained in the above two cases, together with the experimental data from the LHCb [7, 12], Belle [13] and BaBar [108] collaborations. For the NP contributions, we consider only the uncertainties of the model parameters. One can see clearly that there is only a small impact on P_2 and P'_5 in case A, where the chirality-flipped operators $O'_{9,10}$ are absent, while in case B P'_5 could be increased significantly to be consistent with the experimental data and reduce P_2 when the dimuon invariant mass squared q^2 is higher than the zero-crossing point q_0^2 . Numerical results for the zero-crossing points of P_2 (nonzero one) and P'_5 are given in Table 2, both within the SM and in the A2HDM. It is observed that the impact on q_0^2 in case B is more pronounced

Table 2: The zero-crossing points of P_2 (nonzero one) and P'_5 both within the SM and in the A2HDM.

	SM	Case A	Case B
$q_0^2(P_2)$	$3.43^{+0.33}_{-0.32}$	(3.02, 3.90)	(3.02, 4.79)
$q_0^2(P'_5)$	$2.02^{+0.19}_{-0.15}$	(1.77, 2.32)	(1.79, 4.85)

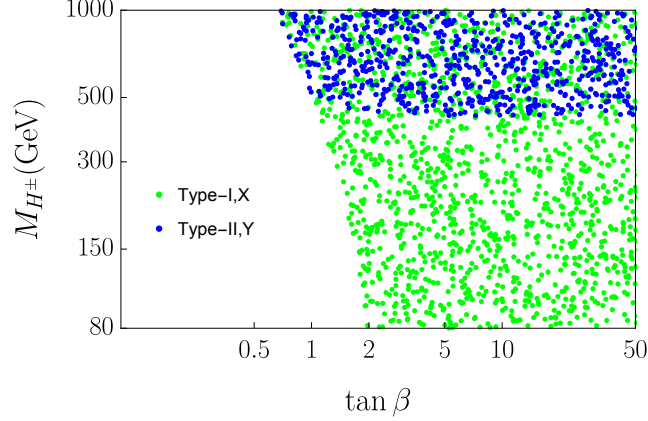


Figure 7: Allowed regions in the $\tan \beta - M_{H^\pm}$ plane corresponding to different \mathcal{Z}_2 -symmetric 2HDMs, under the constraint from Eq. (4.3) as well as the bounds on $C_{9,10}^{H^\pm}$ and $C'_{9,10}{}^{H^\pm}$ from Eqs. (4.4) and (4.5).

than in case A.

4.4 2HDMs with \mathcal{Z}_2 symmetries

In the generic 2HDMs with discrete \mathcal{Z}_2 symmetries, the three alignment parameters ς_f will be reduced to a single parameter $\tan \beta = v_2/v_1 \geq 0$, as indicated in Table 1. There are, therefore, only two model parameters $\tan \beta$ and M_{H^\pm} in the Wilson coefficients $C_{7,9,10}^{(\prime)H^\pm}$. We show in Figure 7 the allowed regions in the $\tan \beta - M_{H^\pm}$ plane corresponding to the four different types of 2HDMs with \mathcal{Z}_2 symmetries. As $C_{7,9,10}^{(\prime)H^\pm}$ do not depend on the parameter ς_ℓ , the type I (II) and type X (Y) models are indistinguishable from each other. However, one can clearly distinguish types I and X from types II and Y models. As shown in Figure 7, the bound $M_{H^\pm} > 432 \text{ GeV}$ is obtained for types II and Y 2HDMs, while there is no further bound found for M_{H^\pm} in types I and X 2HDMs with sizable $\tan \beta$.

In the inert 2HDM [38], on the other hand, $\varsigma_u = \varsigma_d = 0$ leads to $C_{7,9,10}^{(\prime)H^\pm} = 0$, and hence the angular observables remain unaffected.

5 Conclusions

In this paper, we have presented a complete one-loop calculation of the SD Wilson coefficients $C_{7,9,10}^{(\prime)\text{H}^\pm}$ due to the charged-scalar exchanges through the Z^0 - and γ -penguin diagrams within the A2HDM. For $C_{9,10}^{\prime\text{H}^\pm}$, although being suppressed by the factor $\bar{m}_b \bar{m}_s / M_W^2$, they could play an important role in interpreting the observed P_5' anomaly in the decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$, when the model parameter $|\varsigma_d|$ is large.

Under the constraints from the branching ratio $\mathcal{B}(B \rightarrow X_s \gamma)$ and the recent global fit results of $b \rightarrow s \ell^+ \ell^-$ data, we have obtained the allowed parameter spaces in the $\varsigma_u - \varsigma_d$ plane, corresponding to three representative charged-scalar masses. We found that $C_{9,10}^{\text{H}^\pm}$ play a major role in the small $|\varsigma_d|$ region ($|\varsigma_d| < 1$), while $C_{9,10}^{\prime\text{H}^\pm}$ are most important when the model parameter ς_u approaches to zero. When ς_u is far away from zero and $|\varsigma_d| \geq 1$, on the other hand, the impact of $C_7^{\text{H}^\pm}$ will become more significant. Within the constrained parameter space, numerically, the effects of these NP Wilson coefficients can be divided into the following two cases:

Case A: $C_{7,9,10}^{\text{H}^\pm}$ are sizable, but $C_{9,10}^{\prime\text{H}^\pm} \simeq 0$;

Case B: $C_7^{\text{H}^\pm}$ and $C_{9,10}^{\prime\text{H}^\pm}$ are sizable, but $C_{9,10}^{\text{H}^\pm} \simeq 0$.

We have then discussed their impacts on the angular observables P_2 and P_5' in the decay $B^0 \rightarrow K^{*0} \mu^+ \mu^-$. It is found that there is only a small impact on P_2 and P_5' in case A, while the case B could obviously increase P_5' to be consistent with the experimental data and reduce P_2 when the dimuon invariant mass squared q^2 is higher than the zero-crossing point.

Finally, we have explored the constraints on $\tan \beta$ and M_{H^\pm} in five types of \mathcal{Z}_2 -symmetric 2HDMs. The role of chirality-flipped operators $O'_{9,10}$ becomes much more important for large values of $\tan \beta$. Even with the current data, the types I and X and types II and Y could be clearly distinguished from each other. In the inert 2HDM, on the other hand, the fact that $\varsigma_u = \varsigma_d = 0$ leads to $C_{7,9,10}^{(\prime)\text{H}^\pm} = 0$, making the angular observables unaffected.

Future precise measurements of the angular observables in $b \rightarrow s \ell^+ \ell^-$ decays, especially with a finer binning of q^2 , would be very helpful to provide a more definite answer concerning the observed anomalies by the LHCb and Belle collaborations, restricting further or even deciphering the NP models.

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A Basic function

The basic functions $F_i(x)$ introduced in Eqs. (3.12) and (3.13) are defined, respectively, as

$$F_0(x) = \ln x, \quad (\text{A.1})$$

$$F_1(x) = \frac{x}{4-4x} + \frac{x \ln x}{4(x-1)^2}, \quad (\text{A.2})$$

$$F_2(x) = \frac{x}{96(x-1)} - \frac{x^2 \ln x}{96(x-1)^2}, \quad (\text{A.3})$$

$$F_3(x) = \frac{x}{8} \left[\frac{x-6}{x-1} + \frac{(3x+2) \ln x}{(x-1)^2} \right], \quad (\text{A.4})$$

$$F_4(x) = -\frac{3x(x-3)}{32(x-1)} + \frac{x(x^2-8x+4) \ln x}{16(x-1)^2}, \quad (\text{A.5})$$

$$F_5(x) = \frac{-19x^3+25x^2}{36(x-1)^3} + \frac{(5x^2-2x-6)x^2 \ln x}{18(x-1)^4}, \quad (\text{A.6})$$

$$F_6(x) = \frac{8x^3+5x^2-7x}{12(x-1)^3} - \frac{(3x-2)x^2 \ln x}{2(x-1)^4}, \quad (\text{A.7})$$

$$F_7(x) = \frac{x(53x^2+8x-37)}{108(x-1)^4} + \frac{x(-3x^3-9x^2+6x+2) \ln x}{18(x-1)^5}, \quad (\text{A.8})$$

$$F_8(x) = \frac{x(18x^4+253x^3-767x^2+853x-417)}{540(x-1)^5} - \frac{x(3x^4-6x^3+3x^2+2x-3) \ln x}{9(x-1)^6}. \quad (\text{A.9})$$

B Wilson coefficients in A2HDM

The coefficients of the different combinations of the couplings ς_u and ς_d in Eqs. (3.14)–(3.19) are given, respectively, by

$$C_{7,\text{uu}} = -\frac{1}{6}F_6(y_t), \quad (\text{B.1})$$

$$C_{7,\text{ud}} = -\frac{4}{3}F_1(y_t) - \frac{80}{17}F_2(y_t) - \frac{3}{17}F_5(y_t) + \frac{1}{17}F_6(y_t), \quad (\text{B.2})$$

$$C_{9,\text{uu}} = \frac{8}{9}F_1(y_t) - \frac{896}{51}F_2(y_t) - \frac{1}{17}F_5(y_t) - \frac{14}{153}F_6(y_t) - \frac{x_t}{2} \left(-4 + \frac{1}{\sin^2 \theta_W} \right) F_1(y_t), \quad (\text{B.3})$$

$$C_{10,\text{uu}} = \frac{x_t}{2 \sin^2 \theta_W} F_1(y_t), \quad (\text{B.4})$$

$$C'_{9,\text{uu}} = \frac{y_t}{x_t} F_8(y_t), \quad (\text{B.5})$$

$$C'_{9,\text{ud}} = \frac{y_t}{x_t} F_7(y_t), \quad (\text{B.6})$$

$$C'_{9,\text{dd}} = \frac{y_t}{x_t} \left[\frac{2}{9}F_0(x_t) + \frac{20}{9}F_1(y_t) + \frac{928}{51}F_2(y_t) - \frac{2}{17}F_5(y_t) - \frac{11}{153}F_6(y_t) \right], \quad (\text{B.7})$$

$$C'_{10,\text{uu}} = -\frac{1}{17} \left[80F_2(y_t) + 3F_5(y_t) - F_6(y_t) \right], \quad (\text{B.8})$$

$$C'_{10,\text{ud}} = \frac{1}{\sin^2 \theta_W} \left[-\frac{1}{12}F_1(y_t) + \frac{30}{17}F_2(y_t) + \frac{9}{136}F_5(y_t) - \frac{3}{136}F_6(y_t) \right] - \frac{1}{6} \left(-4 + \frac{1}{\sin^2 \theta_W} \right) F_1(y_t), \quad (\text{B.9})$$

$$C'_{10,\text{dd}} = -\frac{1}{\sin^2 \theta_W} \left[\frac{1}{2}F_1(y_t) + F_2(y_t) \right] + \left(-4 + \frac{1}{\sin^2 \theta_W} \right) F_2(y_t), \quad (\text{B.10})$$

and for the Wilson coefficient $C_8^{\text{H}\pm}$, we have [81]

$$C_{8,\text{uu}} = \frac{1}{34} \left[720F_2(y_t) + 27F_5(y_t) + 8F_6(y_t) \right], \quad (\text{B.11})$$

$$C_{8,\text{ud}} = 2F_1(y_t) - \frac{1}{17} \left[240F_2(y_t) + 9F_5(y_t) - 3F_6(y_t) \right]. \quad (\text{B.12})$$

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